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On the $1/m^2$ corrections to the form-factors in $\Lambda_b \to \Lambda_c e \bar{\nu}_e$ decays

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Abstract

Model-independent bounds are presented for the form-factors describing the semileptonic decay $\Lambda_b \to \Lambda_c e \bar{\nu}_e$ at the end-point of the lepton energy spectrum. For the axial-vector form-factor $f_1^A(1)$ we obtain $f_1^A(1) \leq 0.960 \div 0.935$, where the uncertainty arises from the value of μ_π^2, the average kinetic energy of the b quark in the Λ_b baryon. A bound is given on the forward matrix element of the axial current in a Λ_b baryon which is relevant for the inclusive semileptonic decays of the polarized Λ_b decays.

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1. Recently it has been shown [1] how to obtain model-independent bounds on form-factors describing semileptonic transitions of heavy hadrons. It had been known for some time that some of the form-factors describing the decay $B \to D^* e\bar{\nu}_e$ do not receive $1/m$ corrections at the so-called "no-recoil" kinematical point [2]. The only corrections to the heavy-quark symmetry limit appear at order $1/m^2$ and they have been evaluated in the past by making use of model-dependent results [4, 5, 6, 7]. In the paper [1] a model-independent bound on the magnitude of these corrections has been given. The aim of this paper is to present a similar bound on the $1/m^2$ corrections to the heavy quark symmetry predictions for the semileptonic decay $\Lambda_b \to \Lambda_c e\bar{\nu}_e$.

2. We begin by describing the kinematics of the decay. The Λ_b baryon is polarized with the spin vector s and is taken to be at rest in the laboratory frame. It decays into a final hadronic state containing one charmed quark of invariant mass m_X and a lepton pair $e\bar{\nu}_e$ of 4-momentum q. The inclusive decay width summed over all final hadronic states with the same invariant mass can be expressed in terms of the hadronic tensor

$$W_{\mu\nu} = (2\pi)^3 \sum_X \langle \Lambda_b(v,s)|J^\dagger_\mu(0)|X\rangle\langle X|J_\nu(0)|\Lambda_b(v,s)\rangle\delta(p_X - p + q) \tag{1}$$

with $p_X^2 = m_X^2$ and $p = m_{\Lambda_b} v$ the momentum of the Λ_b and J is the weak current. We normalize the $|\Lambda_b\rangle$ state as usual to v_0. The most general decomposition into covariants for the hadronic tensor $W_{\mu\nu}$ has the form

$$\begin{aligned} W_{\mu\nu} &= -g_{\mu\nu}W_1 + v_\mu v_\nu W_2 - i\epsilon_{\mu\nu\alpha\beta}v^\alpha q^\beta W_3 + q_\mu q_\nu W_4 + (q_\mu v_\nu + q_\nu v_\mu)W_5 \\ &- q\cdot s\left[-g_{\mu\nu}G_1 + v_\mu v_\nu G_2 - i\epsilon_{\mu\nu\alpha\beta}v^\alpha q^\beta G_3 + q_\mu q_\nu G_4 + (q_\mu v_\nu + q_\nu v_\mu)G_5\right] \\ &+ (s_\mu v_\nu + s_\nu v_\mu)G_6 + (s_\mu q_\nu + s_\nu q_\mu)G_7 + i\epsilon_{\mu\nu\alpha\beta}v^\alpha s^\beta G_8 \\ &+ i\epsilon_{\mu\nu\alpha\beta}q^\alpha s^\beta G_9\,. \end{aligned} \tag{2}$$

Here the same parametrization has been used as in [12].

The hadronic tensor $W_{\mu\nu}$ is given by the discontinuity of the forward scattering amplitude

$$T_{\mu\nu}(q\cdot v, q^2) = -i\int \mathrm{d}^4x e^{-iq\cdot x}\langle\Lambda_b(v,s)|TJ^\dagger_\mu(x)J_\nu(0)|\Lambda_b(v,s)\rangle\,, \tag{3}$$

across the cut in the complex plane of the variable $v\cdot q$ which corresponds to the process under consideration. In our case the cut extends from $\sqrt{q^2}$ to $1/(2m_{\Lambda_b})(m^2_{\Lambda_b} - m^2_{\Lambda_c} + q^2)$. The invariant mass of the final hadronic state varies from m_{Λ_c} on the right edge of the cut up to $m_{\Lambda_b} - \sqrt{q^2}$ on the left edge.

We will parametrize the forward scattering amplitude $T_{\mu\nu}$ in terms of invariant form-factors $T_{1,5}$ and $S_{1,9}$ (defined in analogy to $W_{1,5}$ and $G_{1,9}$ in (2)).

For our purposes it will prove more convenient to study the analyticity properties of $T_{\mu\nu}$ as a function of complex $v \cdot q$ at fixed velocity v' of the final hadronic state (instead of q^2). In these variables, for a given $\omega = v \cdot v'$, the cut corresponding to the physical decay process ranges between

$$m_{\Lambda_b} \frac{\sqrt{\omega^2 - 1}}{\omega + \sqrt{\omega^2 - 1}} \leq v \cdot q \leq m_{\Lambda_b} - m_{\Lambda_c} \omega \,. \tag{4}$$

The invariant mass of the final hadronic state correspondingly takes values between

$$m_{\Lambda_b} \frac{1}{\omega + \sqrt{\omega^2 - 1}} \geq m_X \geq m_{\Lambda_c} \tag{5}$$

and q^2 is a function of $v \cdot q$ along the cut, given by

$$q^2 = -m_{\Lambda_b}(m_{\Lambda_b} - 2v \cdot q) + \frac{1}{\omega^2}(m_{\Lambda_b} - v \cdot q)^2 \,. \tag{6}$$

The discontinuity of $T_{\mu\nu}$ across this cut is related to the hadronic tensor (1) by

$$W_{\mu\nu} = -\frac{1}{2\pi i} \mathrm{disc}\, T_{\mu\nu} \,. \tag{7}$$

In the following we will construct two sum rules for the form-factors describing the semileptonic decays of the Λ_b baryon in analogy to the ones presented in [1]. The method is based on the positivity of certain combinations of invariant form-factors in (2). These can be obtained from inequalities of the form

$$n^\mu n^{*\nu} W_{\mu\nu} \geq 0 \tag{8}$$

and follow from the fact that this quantity is directly related to the decay rate $\Lambda_b \to W(q, n_\mu) X_c$ of the Λ_b baryon into a virtual W boson with the polarization vector n. Let us work in the rest frame of the Λ_b baryon.

One positive-definite combination of invariant form-factors in (2) is obtained by using in (8) a longitudinal polarization vector. This gives

$$W_L = W_1 + \frac{q_0^2 - q^2}{q^2} W_2 - q \cdot s \left(G_1 + \frac{q_0^2 - q^2}{q^2} G_2 \right) + 2 \frac{q_0}{q^2} (s \cdot q) G_6 \,. \tag{9}$$

Another possible choice is the total decay rate into a circularly polarized W boson, which is proportional to

$$W_{T_{L,R}} = W_1 - q \cdot s G_1 \pm \left(\sqrt{q_0^2 - q^2}(W_3 - q \cdot s G_3) + \frac{q \cdot s}{\sqrt{q_0^2 - q^2}} (G_8 + q_0 G_9) \right) \,. \tag{10}$$

The total decay rate into a transversally polarized W is

$$W_T = 2W_1 - 2q\cdot sG_1\,. \tag{11}$$

Finally, the case of a scalar (or "temporal") polarization gives

$$\begin{aligned} W_0 &= -W_1 + \frac{q_0^2}{q^2}W_2 + q^2W_4 + 2q_0W_5 \\ &- q\cdot s\left(-G_1 + \frac{q_0^2}{q^2}G_2 + q^2G_4 + 2q_0G_5\right) + 2\frac{q_0}{q^2}(s\cdot q)G_6 + 2(s\cdot q)G_7\,. \end{aligned} \tag{12}$$

These quantities must be positive for any possible orientation of the spin vector $\vec{s}$ with respect to $\vec{q}$ and for any current J in (1).

3. The forward scattering amplitude $T_{\mu\nu}$ can be given a representation in terms of physical intermediate states as

$$\begin{aligned} T_{\mu\nu} &= \sum_{X(p_X=m_{\Lambda_b}v-q)} \frac{\langle\Lambda_b|J_\mu^\dagger(0)|X\rangle\langle X|J_\nu(0)|\Lambda_b\rangle}{m_{\Lambda_b} - E_X - q_0 + i\epsilon} \\ &+ \sum_{X(p_X=m_{\Lambda_b}v+q)} \frac{\langle\Lambda_b|J_\nu(0)|X\rangle\langle X|J_\mu^\dagger(0)|\Lambda_b\rangle}{m_{\Lambda_b} - E_X + q_0 + i\epsilon}\,. \end{aligned} \tag{13}$$

Only the first term above gives a contribution to the discontinuity across the cut corresponding to the decay process of interest. We will be mainly interested in the contribution of the intermediate state Λ_c. We will consider first the case of an axial current $J_\mu = \bar{c}\gamma_\mu\gamma_5 b$. The corresponding matrix element of the axial current in (13) can be parametrized as usual in terms of 3 form-factors, defined through

$$\langle\Lambda_c(v',s')|\bar{c}\gamma_\mu\gamma_5 b|\Lambda_b(v,s)\rangle = \bar{u}(v',s')[f_1^A\gamma_\mu + f_2^Av_\mu + f_3^Av'_\mu]\gamma_5 u(v,s)\,. \tag{14}$$

We have used a parametrization appropriate to the heavy-mass limit, in terms of the velocities of the particles involved. The spinors are normalized so that $\bar{u}u = 1$.

By taking the imaginary part of (13), the hadronic tensor $W_{\mu\nu}$ can be also written as a sum over intermediate states. The contribution of the lowest-lying state Λ_c to the positive-definite combinations of structure functions defined in (9,10,12) is

$$W_L = \frac{\omega+1}{2q^2}\left((m_{\Lambda_b} - m_{\Lambda_c})f_1^A - m_{\Lambda_c}(\omega-1)f_2^A - m_{\Lambda_b}(\omega-1)f_3^A\right)^2 \tag{15}$$

$$W_{T_{L,R}} = |f_1^A|^2\frac{1+\omega}{2}(1\pm\hat{n}\cdot\vec{s}) \tag{16}$$

$$W_0 = \frac{\omega-1}{2q^2}\left((m_{\Lambda_b} + m_{\Lambda_c})f_1^A - (m_{\Lambda_b} - m_{\Lambda_c}\omega)f_2^A - (m_{\Lambda_b}\omega - m_{\Lambda_c})f_3^A\right)^2 \tag{17}$$

A factor of $1/\omega\delta(q_0 - m_{\Lambda_b} + m_{\Lambda_c}\omega)$ has been removed from these expressions. Here $\hat{n}$ is a unit vector collinear with the vector $\vec{q}$. If $|\vec{q}| = 0$, it defines the quantization axis along which the spin of the virtual W boson is aligned. The above formulas can be simply related to helicity amplitudes for the process $\Lambda_b(\lambda_{\Lambda_b}) \to \Lambda_c(\lambda_{\Lambda_c}) + W(\lambda_W)$ listed for different other cases of physical interest in [8].

For the case of a vector current $J_\mu = \bar{c}\gamma_\mu b$, we define in an analogous way to (14)

$$\langle \Lambda_c(v', s')|\bar{c}\gamma_\mu b|\Lambda_b(v, s)\rangle = \bar{u}(v', s')[f_1^V \gamma_\mu + f_2^V v_\mu + f_3^V v'_\mu]u(v, s)\,. \tag{18}$$

In terms of these form-factors, the contribution of the intermediate state Λ_c reads

$$W_L = \frac{\omega - 1}{2q^2}\Big((m_{\Lambda_b} + m_{\Lambda_c})f_1^V + m_{\Lambda_c}(\omega + 1)f_2^V + m_{\Lambda_b}(\omega + 1)f_3^V\Big)^2 \tag{19}$$

$$W_{T_{L,R}} = |f_1^V|^2 \frac{\omega - 1}{2}(1 \pm \hat{n}\cdot\vec{s}) \tag{20}$$

$$W_0 = \frac{\omega + 1}{2q^2}\Big((m_{\Lambda_b} - m_{\Lambda_c})f_1^V + (m_{\Lambda_b} - m_{\Lambda_c}\omega)f_2^V + (m_{\Lambda_b}\omega - m_{\Lambda_c})f_3^V\Big)^2 \tag{21}$$

As before, a factor of $1/\omega\delta(q_0 - m_{\Lambda_b} + m_{\Lambda_c}\omega)$ has to be added on the r.h.s..

4. On the other hand, it has been recently shown [9, 10, 11, 12, 13] that it is possible to reliably calculate the forward scattering matrix element (3) in a region which is far away from the physical cuts. The result takes the form of an expansion in powers of $1/m_{c,b}$ and depends on a few nonperturbative matrix elements of dimension-6 operators taken between heavy hadron states. For the case of an axial (vector) current $J_\mu = \bar{c}\gamma_\mu\gamma_5 b$, the results for $T_{1,5}$ can be taken from [11]. As for the spin-dependent structure functions one can see from (9,10,12) that at the equal-velocity point $\omega = 1$, only $G_{8,9}$ contribute. The corresponding amplitudes $S_{8,9}$ can be straightforwardly calculated following the method given in [11, 12] with the result (for $J_\mu = \bar{c}\gamma_\mu\gamma_5 b$; the result for the vector current case can be obtained by making the replacement $m_c \to -m_c$)

$$\begin{aligned} S_8 &= \frac{1}{\Delta}\left(1 + \frac{\mu_s^2}{m_b^2} + \frac{\mu_\pi^2}{6m_b^2}\right)(m_b + m_c) + \frac{5\mu_\pi^2}{3m_b\Delta^2} v\cdot q(m_b + m_c) \\ &+ \frac{4\mu_\pi^2}{3\Delta^3}[(v\cdot q)^2 - q^2](m_b + m_c) \end{aligned} \tag{22}$$

$$\begin{aligned} S_9 &= -\frac{1}{\Delta}\left(1 + \frac{\mu_s^2}{m_b^2}\right) - \frac{\mu_\pi^2}{m_b\Delta^2}\left(v\cdot q + \frac{2}{3}(m_b + m_c)\right) \\ &- \frac{4\mu_\pi^2}{3\Delta^3}[(v\cdot q)^2 - q^2] \end{aligned} \tag{23}$$

where

$$\Delta = m_b^2 - 2m_b q_0 + q^2 - m_c^2 + i\epsilon \tag{24}$$

and

$$\mu_\pi^2 = -\langle \Lambda_b(v,s)|\bar{h}_b(iD)^2 h_b|\Lambda_b(v,s)\rangle_{spin-averaged}\,. \tag{25}$$

μ_π^2 represents the average kinetic energy of the b quark inside a Λ_b baryon. At present its numerical value is only poorly known. It has been shown, on quite general grounds, that it must be positive [15]. A previous estimate [4] of the corrections of order $1/m^2$ to the $\Lambda_b \to \Lambda_c$ form-factors used $\mu_\pi^2 \simeq -1$ GeV^2. The corresponding parameter in the $B - B^*$ system has the approximate value 0.5 GeV^2. Using the mass formula

$$m_{\Lambda_b} = m_b + \bar{\Lambda} + \frac{\mu_\pi^2}{2m_b} \tag{26}$$

and a similar one for the Λ_c baryon, together with the mass values [14] $m_{\Lambda_b} = 5.641$ GeV, $m_{\Lambda_c} = 2.285$ GeV, $m_b = 4.8$ GeV and $m_c = 1.39$ GeV, yields

$$\mu_\pi^2 = 0.211\,\mathrm{GeV}^2\,. \tag{27}$$

For the purpose of illustration we will vary the value of μ_π^2 between 0.211 and 0.500 GeV^2.

The other parameter in (22,23) is μ_s^2, which is defined as the $1/m_b^2$-correction to the forward matrix element of the axial current in a Λ_b state

$$\langle \Lambda_b(v,s)|\bar{b}\gamma_\mu\gamma_5 b|\Lambda_b(v,s)\rangle = \left(1 + \frac{\mu_s^2}{m_b^2} + \cdots\right)\bar{u}\gamma_\mu\gamma_5 u\,. \tag{28}$$

This matrix element does not receive any corrections at order $1/m_b$ [3]. μ_s^2 is related to the parameter ϵ_b defined in [12] by $\mu_s^2 = m_b^2\epsilon_b$. Very little is known about its precise numerical value. We will show below that it must be negative and will give an upper bound for its value.

Inserting q^2 from (6) in the theoretical expression for $T_{\mu\nu}$ and taking the imaginary part yields the QCD prediction for the hadronic tensor, which is a singular function of q_0 consisting of δ functions and its derivatives. Integrating the QCD prediction for W_1 over the physical cut in the q_0 complex plane at fixed $\omega = 1$ one obtains

$$\int \mathrm{d}q_0 W_1(q_0, \omega = 1) = 1 - \frac{\mu_\pi^2}{6m_c m_b} - \frac{\mu_\pi^2}{4m_b^2} - \frac{\mu_\pi^2}{4m_c^2}\,. \tag{29}$$

As one can see from (11), this quantity is positive definite, which means that it represents an upper bound for the contribution of any physical state to the hadronic

tensor. In particular, from (15,16) one obtains a bound on the invariant form-factor $f_1^A(1)$ which is relevant for the decay rate of the process $\Lambda_b \to \Lambda_c e\bar{\nu}_e$ near the endpoint of the electron spectrum:

$$|f_1^A(1)|^2 \leq 1 - \frac{\mu_\pi^2}{6m_c m_b} - \frac{\mu_\pi^2}{4m_b^2} - \frac{\mu_\pi^2}{4m_c^2}\,. \tag{30}$$

This bound is similar to the bound recently derived in [1]. Numerically, this bound is

$$1 - |f_1^A(1)|^2 \geq 0.165\mu_\pi^2\,(\mathrm{GeV}^2) = (3.5 \div 8.3)\%\,. \tag{31}$$

This result is comparable with the estimate of the same quantity in [4] which gave about 6.2%.

This bound can be improved if the contribution of the spin-dependent structure functions is taken into account. Let us consider for example the combination W_{T_L} defined in (10), at $\omega = 1$. The Λ_c contribution to its integral along the cut is

$$\int \mathrm{d}q_0 W_{T_L}(q_0, \omega = 1) = |f_1^A(1)|^2(1 + \hat{n}\cdot\vec{s}\,) \tag{32}$$

whereas the QCD prediction for the same quantity is

$$\begin{aligned} \int \mathrm{d}q_0 W_{T_L}(q_0, \omega = 1) \;\; &= \;\; \left(1 - \frac{\mu_\pi^2}{6m_c m_b} - \frac{\mu_\pi^2}{4m_b^2} - \frac{\mu_\pi^2}{4m_c^2}\right)(1 + \hat{n}\cdot\vec{s}\,) \\ &+ \;\; (n\cdot\vec{s}\,)\frac{1}{m_b^2}(\mu_s^2 + \frac{\mu_\pi^2}{3})\,. \end{aligned} \tag{33}$$

Requiring the positivity of this quantity for any value of $\hat{n}\cdot\vec{s}$ gives the inequality

$$\mu_s^2 + \frac{\mu_\pi^2}{3} \leq 0\,, \tag{34}$$

from which one obtains an upper bound on μ_s^2

$$\mu_s^2 \leq -\frac{\mu_\pi^2}{3} \simeq (-0.070 \div -0.167)\,\mathrm{GeV}^2\,. \tag{35}$$

In [4] μ_s^2 has been estimated to be $\mu_s^2 = -\mu_\pi^2/3$, assuming that the contribution of terms arising from double insertions of chromomagnetic operator can be neglected. The present method gives also the sign of the corrections to this estimate.

The inequality (34) implies that the bound (30) on $|f_1^A(1)|$ can be improved, provided that a determination (calculation) of μ_s^2 becomes available. For the moment, we will restrict ourselves to the simple bound (30).

In a completely analogous way one can obtain a different bound on the vector-current form-factors $f_{1,2,3}^V$ defined in (18). This time the bound results from requiring

the positivity of the scalar combination of form-factors (12). From (21), the contribution of the Λ_c state reads

$$\int \mathrm{d}q_0 W_0(q_0, \omega = 1) = (f_1^V(1) + f_2^V(1) + f_3^V(1))^2 \,, \tag{36}$$

whereas the QCD analysis yields for the same quantity

$$\int \mathrm{d}q_0 W_0(q_0, \omega = 1) = 1 - \frac{\mu_\pi^2}{4}\left(\frac{1}{m_c} - \frac{1}{m_b}\right)^2 \,. \tag{37}$$

It has been shown [3] that the combination of form-factors appearing in (36) receives no $1/m$ corrections at $\omega = 1$. Putting together (36) and (37) gives a lower bound on the magnitude of the $1/m^2$ corrections to the leading-order result:

$$1 - (f_1^V(1) + f_2^V(1) + f_3^V(1))^2 \geq \frac{\mu_\pi^2}{4}\left(\frac{1}{m_c} - \frac{1}{m_b}\right)^2 = 0.065\mu_\pi^2 \simeq (1.4 \div 3.2)\% \,. \tag{38}$$

which is in agreement with the 19% estimate of this quantity in [4]. On the other hand, in [7] $\sum f_i^V(1)$ was found to be larger than unity. As will be shown below, the radiative corrections could account for such an increase.

These bounds get changed when radiative corrections are incorporated. Their effect is to modify the heavy quark prediction for the form-factors at leading order in $1/m$ by a multiplicative factor. To one-loop order, the factor multiplying $f_1^A(1)$ is [16]

$$\eta_A^{pert} = 1 + \frac{\alpha_s}{\pi}\left(\frac{m_b + m_c}{m_b - m_c}\log\frac{m_b}{m_c} - \frac{8}{3}\right) \simeq 0.958 \,, \tag{39}$$

and the factor multiplying the form-factor combination $f_1^V(1) + f_2^V(1) + f_3^V(1)$

$$\eta_V^{pert} = 1 + \frac{\alpha_s}{\pi}\left(\frac{m_b + m_c}{m_b - m_c}\log\frac{m_b}{m_c} - 2\right) \simeq 1.025 \tag{40}$$

where $\Lambda_{QCD} = 0.250$ GeV has been used. As a result we finally obtain

$$f_1^A(1) \leq 0.960 \div 0.935 \tag{41}$$

$$f_1^V(1) + f_2^V(1) + f_3^V(1) \leq 1.005 \div 0.996 \,. \tag{42}$$

In principle these bounds could be improved if the contributions of any other states to the integrals of the positive-definite quantities (11,12) are known or if they can be estimated in some way. In particular, the contribution of two-body intermediate states like $\Lambda_c\eta$ and $\Sigma_c\pi$ can be estimated by making use of the Heavy Hadron Chiral Perturbation Theory [17, 18]. Unfortunately, the coupling of the light Goldstone bosons to the antitriplet of heavy baryons containing one heavy quark is at present unknown and any estimate of these contributions will necessarily entail model-dependent assumptions.

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